

Kaon Photoproduction in a Confining and Covariant Diquark Model*

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Abstract

Baryons are modeled as bound states of scalar or axialvector diquarks and a constituent quark which interact through quark exchange. This description results as an approximation to the relativistic Faddeev equation for three quarks. The corresponding effective Bethe-Salpeter equation is solved, and fully four-dimensional wave functions for both octet and decuplet baryons are obtained. Confinement of quarks and diquarks is incorporated by modifying their propagators. Results are given for kaon photoproduction which is one of the applications currently under investigation in this model.

Motivation

In this talk it will be argued that in the study of baryonic structure kaon photoproduction is an especially interesting hadronic reaction. Photo- and electroproduction of pseudoscalar mesons have been studied intensively in isobar resonance models [1,2] and coupled channel calculations of hadronic interactions [3]. These models are capable of describing the wealth of measured data reasonably well. However, by construction they do not allow for an interpretation of the production process in terms of baryon substructure. The aim of the investigation reported here is to clarify whether the notion of diquarks as effective constituents of baryons in meson production close to threshold is helpful in the understanding of subnucleonic physics.

Approaches to the substructure of baryons include nonrelativistic quark models, various sorts of bag models and different types of solitons [4]. Most of these models are designed to work in the low energy region and generally do not match the calculations within perturbative QCD. During this conference impressive reports on the great experimental progress in medium energy physics have been given. This underlines the high demand for models describing baryon physics in this region. Therefore we propose a fully covariant description of baryon structure in the framework of a diquark-quark-model of baryons.

Our motivation to choose such an approach is based on two sources. On the one hand, when starting with the fully relativistic Faddeev equation for bound states of three quarks, diquarks appear as effective degrees of freedom. These diquarks stand for (potentially weakly) correlated quark-quark pairs inside baryons. On the other hand, diquarks as constituents of baryons are naturally obtained when one starts with an NJL-type model of colour octet flavour singlet quark currents [5]. Although in the limit $N_c \rightarrow \infty$ baryons emerge as solitons of meson fields [6], it can be shown for the case of three colours that both effects, binding through quark exchange in the diquark-quark picture and through mesonic effects, contribute equally [7].

The four-dimensional Bethe-Salpeter equation: masses and wave functions

As stated above, in a first step we want to reduce the complexity of the full three-body problem of the relativistic Faddeev equation for baryons. This can be achieved by approximating the two-quark irreducible T -matrix by separable contributions that can be viewed as loosely bound diquarks. The three-body problem then becomes an effective two-body one, in which bound states appear as the solution of a homogeneous Bethe-Salpeter equation. The attractive interaction between quark and diquark is hereby provided by

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Table 1. Components of the octet baryon wave function with their respective spin and orbital angular momentum. $(\gamma_5 C)$ corresponds to scalar and $(\gamma^\mu C)$, $\mu = 1 \dots 4$, to axialvector diquark correlations. Note that the partial waves in the first row possess a non-relativistic limit. See [8] for further details.

“non-relat.”				
partial waves	$\begin{pmatrix} \chi \\ 0 \end{pmatrix}_{(\gamma_5 C)}$	$\hat{P}^4 \begin{pmatrix} 0 \\ \chi \end{pmatrix}_{(\gamma^4 C)}$	$\begin{pmatrix} i\sigma^i \chi \\ 0 \end{pmatrix}_{(\gamma^i C)}$	$\begin{pmatrix} i\left(\hat{p}^i(\vec{\sigma}\vec{p}) - \frac{\sigma^i}{3}\right)\chi \\ 0 \end{pmatrix}_{(\gamma^i C)}$
spin	1/2	1/2	1/2	3/2
orb.ang.mom.	s	s	s	d
<hr/>				
“relat.”				
partial waves	$\begin{pmatrix} 0 \\ \vec{\sigma}\vec{p}\chi \end{pmatrix}_{(\gamma_5 C)}$	$\hat{P}^4 \begin{pmatrix} (\vec{\sigma}\vec{p})\chi \\ 0 \end{pmatrix}_{(\gamma^4 C)}$	$\begin{pmatrix} 0 \\ i\sigma^i(\vec{\sigma}\vec{p})\chi \end{pmatrix}_{(\gamma^i C)}$	$\begin{pmatrix} 0 \\ i\left(p^i - \frac{\sigma^i(\vec{\sigma}\vec{p})}{3}\right)\chi \end{pmatrix}_{(\gamma^i C)}$
spin	1/2	1/2	1/2	3/2
orb.ang.mom.	p	p	p	p

quark exchange. This interaction is but the minimal correlation needed to reconstitute the Pauli principle. Due to antisymmetry in the color indices and the related symmetrization of all other quantum numbers the Pauli principle leads to an attractive interaction in contrast to “Pauli repulsion” known in conventional few-fermion systems. All unknown and probably very complicated gluonic interactions between two quarks are effectively treated via the parameterization of the diquark propagator and the diquark–quark–quark vertex function. In Ref. [8] we have formalized this procedure by an effective Lagrangian containing constituent quark, scalar diquark and axialvector diquark fields. This leads to a coupled set of Bethe–Salpeter equations for octet and decuplet baryons.

We avoid unphysical thresholds by an effective parameterization of confinement in the quark and diquark propagators. We solve the four-dimensional equations in ladder approximation and obtain wave functions for the octet and decuplet baryons [8]. The Lorentz invariance of our model has been checked explicitly by choosing different frames.

The implementation of the appropriate Dirac and Lorentz representations of the quark and diquark parts of the wave functions leads to a unique decomposition in the rest frame of the baryon. Besides the well known s -wave and d -wave components of non-relativistic formulations of the baryon octet we additionally obtain non-negligible p -wave contributions which demonstrates again the need for covariantly constructed models. Table 1 summarizes the structure of the octet wave function. Each of the eight components is to be multiplied with a scalar function which is given in terms of an expansion in hyperspherical harmonics and is computed numerically.

In order to obtain the mass spectra for the octet and decuplet baryons we explicitly break $SU(3)$ flavour symmetry by a higher strange quark constituent mass. Using the nucleon and the delta mass as input our calculated mass spectra [8] are in good agreement with the experimental ones, see Table 2. The wave functions for baryons with distinct strangeness content but same spin differ mostly due to flavour Clebsch-Gordan coefficients, the respective invariant functions being very similar. Due to its special role among the other baryons, we investigated the Λ hyperon in more detail and discussed its vertex amplitudes. In our approach, the Λ acquires a small flavour singlet admixture which is absent in $SU(6)$

Table 2. Octet and decuplet masses obtained with two different parameter sets. Set I represents a calculation with weakly confining propagators, Set II with strongly confining propagators, see [8]. All masses are given in GeV.

	m_u	m_s	M_Λ	M_Σ	M_Ξ	M_{Σ^*}	M_{Ξ^*}	M_Ω
Set I	0.5	0.65	1.123	1.134	1.307	1.373	1.545	1.692
Set II	0.5	0.63	1.133	1.140	1.319	1.380	1.516	1.665
Exp.			1.116	1.193	1.315	1.384	1.530	1.672



Figure. 1. Typical diagrams to be computed in the diquark spectator picture. Left: kaon photoproduction. Right: strangeness production in proton–proton collisions.

symmetric non-relativistic quark models.

Form Factors

A significant test of our model is the calculation of various form factors [9,10,11]*. The most important ingredient are the fully four-dimensional wave functions described above. It turns out that already the electromagnetic form factors of the nucleon provide severe restrictions for the parameters of the model.

For the pion–nucleon form factor at the soft point, $g_{\pi NN}$, we find good agreement with experiment. For spacelike momenta this form factor falls like a monopole with a large cutoff similar to the behaviour in One-Boson-Exchange (OBE) models [11]. Compared with a calculation including only scalar diquarks [9] we find a lower value for the pion–nucleon coupling at the soft point. Serving as a central ingredient for strangeness production processes the kaon–nucleon–lambda form factor $g_{K\Lambda N}$ is a quantity of special interest. Due to flavour algebra the isospin configuration of the Λ singles out the scalar diquark as the only diquark contributing to nucleon–lambda transitions. Since a pseudoscalar kaon does not couple to the scalar diquark we find ourselves in a comfortable position to handle such transitions. We find that the absolute value of the kaon–nucleon–lambda form factor is always smaller than the pion–nucleon form factor [11]. This can be understood from the facts that the axialvector diquark does not contribute and that the kaon decay constant is larger than the pion decay constant.

Kaon Photoproduction

The SAPHIR collaboration has measured the cross section and the asymmetries for the process $p\gamma \rightarrow K\Lambda$ with high precision. These measurements continue to prompt corresponding calculations within various models, among these are the isobaric model [1,2], the coupled channel approach [3] and models which describe those processes in terms of the dynamics of the baryon substructure [14].

The reaction $p\gamma \rightarrow K\Lambda$ lends itself to a description within the diquark–quark picture. This is for two reasons which greatly simplify the description: first, the scalar diquark is the sole overlap between the wave function of the proton and the lambda, which implies that axialvector diquarks cannot participate in the reaction, and second, the kaon does not couple to the diquark. This leads to the conclusion that the diagram shown in the left half of Fig. 1 and the corresponding crossed diagram are the dominant contributions to kaon photoproduction.

The total cross section for $p\gamma \rightarrow K\Lambda$ is shown in the left panel of Fig. 2. For $E < 1.5$ GeV the data are reproduced quite nicely, whereas our results do not fall off fast enough for higher energies. We found that the total cross section depends very little on the details of the baryon wave functions, however, we observed a rather pronounced dependence on the way confinement is parameterized into the propagators. A closer analysis reveals that production processes like $p\gamma \rightarrow K\Lambda$ probe the propagators of the constituents in the time-like region, i.e. for timelike momenta of the constituents. This is a highly welcome feature

*For a similar calculation of the electromagnetic nucleon form factors within a slightly different diquark model see Ref. [12]

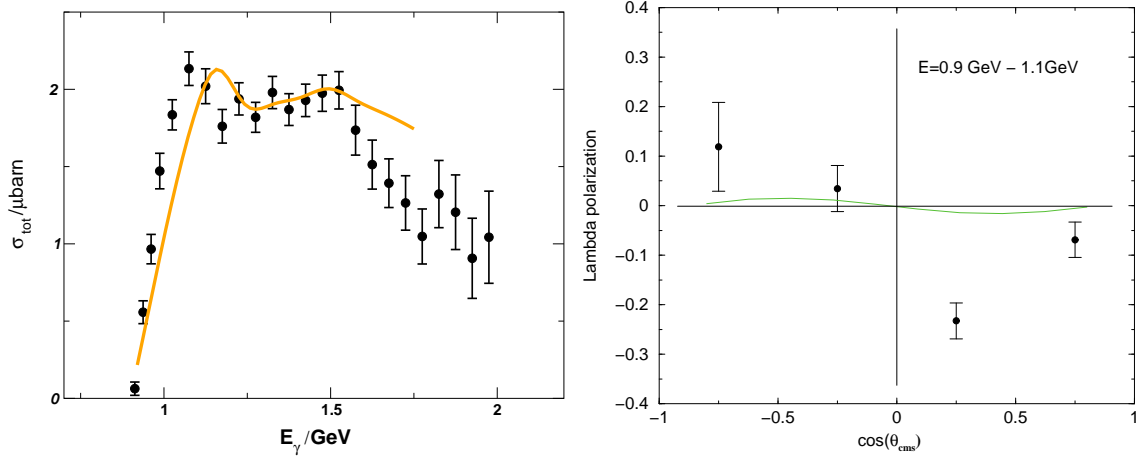


Figure. 2. Total cross section (left panel) and Lambda polarization (right panel) for the reaction $p\gamma \rightarrow K\Lambda$. The data are taken from [15] and are shown together with our results.

which may be used in conjunction with Dyson–Schwinger studies. This approach is tied to the dynamics of QCD and thus gives access to the nonperturbative quark propagator, which enters as a key ingredient in the description of baryons as bound states of quarks and diquarks. Since the Dyson-Schwinger approach gives the propagators for spacelike momenta only, one should take on board the idea, that production processes like $p\gamma \rightarrow K\Lambda$ firmly constrain the extrapolation to timelike momenta.

The lambda polarization is shown in the right panel of Fig. 2 for the energy range 0.9 – 1.1 GeV. Our results fall short of the experimental values, however, the change in sign is reproduced. Comparing with, e.g. [1,2,3] we conclude that apparently none of the current models is able to reproduce the asymmetries to reasonable accuracy.

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